Low-scale inflation in a model of dark energy and dark matter

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We present a complete particle physics model that explains three major problems of modern cosmology: inflation, dark matter and dark energy, and also gives a mechanism for leptogenesis. The model has a new gauge group $SU(2)_Z$ that grows strong at a scale $\Lambda \sim 10^{-3}$ eV. We focus on the inflationary aspects of the model. Inflation occurs with a Coleman-Weinberg potential at a low scale, down to $\sim 6 \times 10^5$ GeV, being compatible with observational data.

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I. INTRODUCTION

The recent three-year Wilkinson Microwave Anisotropy Probe (WMAP3) results [1] have put quite a severe constraint on inflationary models. In particular, new results on the value of the spectral index $n_s = 0.95 \pm 0.02$ are sufficiently "precise" as to rule out many models with an exact Harrison-Zel'dovich-Peebles scale-invariant spectrum with $n_s = 1$ and for which the tensor-to-scalar ratio $r \ll 1$. Any model purported to describe the early inflationary era will have to take into account this constraint. However, by itself, it is not sufficient to narrow down the various candidate models of inflation. In particular, "low-scale" inflationary models are by no means ruled out by the new data. By "low-scale" we refer to models in which the scale that characterizes inflation is several orders of magnitude smaller than a typical Grand Unified Theory (GUT) scale $\sim 10^{15} - 10^{16}$ GeV. It is in this context that we wish to present a model of low-scale inflation which could also describe the dark energy and dark matter [2], [3].

The model of dark energy and dark matter described in [3] involves a new gauge group $SU(2)_Z$ which grows strong at a scale $\Lambda_Z \sim 3\times10^{-3}$ eV starting with the value of the gauge coupling at $\sim 10^{16}$ GeV which is not too different from the Standard Model (SM) couplings at a similar scale. (This is nicely seen when we embed $SU(2)_Z$ and the SM in the unified gauge group E_6 [4].) The model of [3] contains, in addition to the usual SM content, particles which are SM singlets but $SU(2)_Z$ triplets, $\psi^{(Z)}_{(L,R),i}$ with $i = 1, 2$, particles which carry quantum numbers of both gauge groups, $\tilde{\varphi}^{(Z)}_{1,2}$, which are the so-called messenger fields with the decay of $\tilde{\varphi}^{(Z)}_1$ being the source of SM leptogenesis [5], and a singlet complex scalar field,

$$\phi_Z = (\sigma_Z + v_Z)\, \exp(i a_Z/v_Z)\,, \tag{1}$$

whose angular part a_Z is the axion-like scalar. We have defined the radial part of ϕ_Z as the sum of a field σ_Z and a vacuum-expectation-value (v.e.v.) v_Z. The $SU(2)_Z$ instanton-induced potential for a_Z (with two degenerate vacua) along with a soft-breaking term whose dynamical origin is discussed in [6], is one that is proposed in [3] as a model for dark energy. In that scenario, the present universe is assumed to be trapped in a false vacuum of the a_Z potential with an energy density $\sim \Lambda_Z^4$. The exit time to the true vacuum was estimated in [3] and was found to be enormous, meaning that our universe will eventually enter a late inflationary stage.

What might be interesting is the possibility that the real part of ϕ_Z, namely σ_Z, could play the role of the inflaton while the imaginary part, a_Z, plays the role of the "acceleron" as we have mentioned above. This unified description is attractive for the simple reason that *one complex* field describes both phenomena: Early and Late inflation. (This scenario has been exploited earlier [7], [8] in the context of GUT scale inflation.) Although the structure of the potential describing the accelerating universe is determined, in the model of [3], by instanton dynamics of the $SU(2)_Z$ gauge interactions [3], the potential for σ_Z, which would describe the early inflationary universe, is arbitrary as with scalar potentials in general. In this case, the only constraint comes from the requirement that this potential should be of the type that gives the desired spectral index and the right amount of inflation corresponding to the characteristic scale of the model.

We now briefly describe the model of [3]. The key ingredient of that model is the postulate of a new, unbroken gauge group $SU(2)_Z$ which grows strong at a scale $\Lambda_Z \sim 3 \times 10^{-3}$ eV. The model also contains a global symmetry $U(1)^{(Z)}_A$ which is spontaneously broken by the v.e.v. of ϕ_Z, namely $\langle\phi_Z\rangle = v_Z$, and is also *explicitly* broken at a scale $\Lambda_Z \ll v_Z$ by the $SU(2)_Z$ gauge anomaly. Because of this, the pseudo-Nambu-Goldstone boson (PNGB) a_Z acquires a tiny mass as discussed in [3]. Its $SU(2)_Z$ instanton-induced potential used in the false vacuum scenario for the dark energy is given by

$$V_{tot}(a_Z,T) = \Lambda_Z^4 \left[1-\kappa(T)\cos\frac{a_Z}{v_Z}\right] + \kappa(T)\Lambda_Z^4\,\frac{a_Z}{2\pi\, v_Z}\,, \tag{2}$$

where $\kappa(T)=1$ at $T<\Lambda_Z$. ($SU(2)_Z$ instanton effects become important when $\alpha_Z = g_Z^2/4\,\pi \sim 1$ at $\Lambda_Z \sim 3\times 10^{-3}\,eV$.) The universe is assumed to be presently trapped in the false vacuum at $a_Z = 2\pi v_Z$ with an energy density $\sim (3\times 10^{-3}\,eV)^4$. As such, this model mimicks the ΛCDM scenario with $w(a_Z) = \frac{\frac{1}{2}\dot{a}_Z^2 - V(a_Z)}{\frac{1}{2}\dot{a}_Z^2 + V(a_Z)} \approx -1$, at present and for a long time from now on, but not in the distant past [3].

What could be the form of the potential for the *real part*, namely σ_Z, of ϕ_Z? As with any scalar field, the form of the potential is rather arbitrary, with the general constraints being gauge invariance and renormalizability. In this paper, we would like to propose a form of potential for σ_Z which is particularly suited to the discussion of the "low-scale" inflationary scenario: a Coleman-Weinberg (CW) type of potential [9]. (The CW-type of potential has been recently used [10] to describe a GUT-scale inflation using the WMAP3 data.) There are three types of contributions to the potential. The sources of these three types are the following terms in the lagrangian:

a) $\phi_Z - \psi^{(Z)}_{(L,R),i}$ coupling

$$\sum_i K_i\, \bar{\psi}^{(Z)}_{L,i}\, \psi^{(Z)}_{R,i}\, \phi_Z + h.c. \tag{3}$$

Let us recall from [3] that (3) is invariant under the following global $U(1)^{(Z)}_A$ symmetry transformations: $\psi^{(Z)}_{L,i} \to e^{-i\alpha}\,\psi^{(Z)}_{L,i}$, $\psi^{(Z)}_{R,i} \to e^{i\alpha}\,\psi^{(Z)}_{R,i}$, and $\phi_Z \to e^{-2i\alpha}\,\phi_Z$.

b) $\phi_Z - \tilde{\varphi}_1^Z$ mixing (we ignore the $\phi_Z - \tilde{\varphi}_2^Z$ coupling since it is assumed to have a mass of the order of a typical GUT scale)

$$\tilde{\lambda}_{1Z}(\phi_Z^\dagger\,\phi_Z)(\tilde{\varphi}_1^{Z,\dagger}\,\tilde{\varphi}_1^Z) \tag{4}$$

c) σ_Z self-interaction

$$\frac{\lambda}{4!}\sigma_Z^4 \tag{5}$$

Both terms, (4) and (5), arise from the general potential for all fields.

Let us look into constraints on these couplings coming from issues discussed in [3]: dark matter and leptogenesis.

Since the coupling (3) will, in principle, contribute to the CW potential for σ_Z, it is crucial to have an estimate on the magnitude of the Yukawa couplings K_i. In [3], an argument was made as to why it might be possible that $\psi_i^{(Z)}$ could be Cold Dark Matter (CDM) provided

$$m_{\psi_i^{(Z)}} = |K_i| v_Z \le O(200\,GeV)\,, \tag{6}$$

or

$$|K_i| \le O(200\,GeV/v_Z) \tag{7}$$

Roughly speaking, in order for $\Omega_{CDM} \sim O(1)$, the annihilation cross sections for $\psi_i^{(Z)}$ are required to be of the order of weak cross sections. In this case, they are approximately $\sigma_{\text{annihilation}} \sim \alpha_Z^2(m_{\psi_i^{(Z)}})/m^2_{\psi_i^{(Z)}}$ ($\alpha_Z^2(m_{\psi_i^{(Z)}})$ is the coupling evaluated at $E = m_{\psi_i^{(Z)}}$) and have the desired magnitude when $m_{\psi_i^{(Z)}} \sim O(200\,GeV)$ with $\alpha_Z^2(m_{\psi_i^{(Z)}}) \sim \alpha^2_{SU(2)_L}(m_{\psi_i^{(Z)}})$, a characteristic feature of the model of [3].

A second requirement comes from a new mechanism for leptogenesis as briefly mentioned in [3] and described in detail in [5]. This new scenario of leptogenesis involves the decay of a messenger scalar field, $\tilde{\varphi}_1^{(Z)}$, into $\psi_i^{(Z)}$ and a SM lepton. In order to give the correct estimate for the net lepton number, a bound on the mass of $\tilde{\varphi}_1^{(Z)}$ was derived. In [5], it was found that

$$m_{\tilde{\varphi}_1^{(Z)}} \lesssim 1\,TeV\,. \tag{8}$$

This came about when one calculates the interference between the tree-level and one-loop contributions to the decays

$$\tilde{\varphi}_1^{(Z)} \to \bar{\psi}^{(Z)}_{1,2} + l \tag{9}$$

$$\tilde{\varphi}_1^{(Z),*} \to \psi^{(Z)}_{1,2} + \bar{l} \tag{10}$$

where l represents a SM lepton. By requiring that the asymmetry coming from this scenario to be $\epsilon_l^{\tilde{\varphi}_1} \sim -10^{-7}$ in order to obtain the right amount of baryon number asymmetry through the electroweak sphaleron process, [5] came up with the constraint (8) which could be interesting for searches of non-SM scalars at the Large Hadron Collider (LHC). On the other hand, as discussed in [3], the mixing between $\tilde{\varphi}_1^{(Z)}$ and ϕ_Z results in an additional term in the mass squared formula for $\tilde{\varphi}_1^{(Z)}$, namely $2\tilde{\lambda}_{1Z} v_Z^2$. Taking into account the leptogenesis bound (8), one can write

$$\tilde{\lambda}_{1Z} \lesssim (1\,TeV/v_Z)^2\,. \tag{11}$$

All these constraints will be used to estimate the contributions to the effective potential. The $\psi_i^{(Z)}$ fermion loop contribution to the σ_Z CW potential is given by $-|K_i|^4/16\pi^2$, and therefore will be bound by (7). The $\tilde{\varphi}_1^{(Z)}$ loop contribution to the potential is given by $\tilde{\lambda}_{1Z}^2/16\pi^2$ and is constrained by (11). The third contribution c) coming from the σ_Z loop is given by $\lambda^2/16\pi^2$.

There are no constraints on it coming from dark matter or leptogenesis arguments, as we have in the other cases a) and b).

Below, we will constrain both v_Z and the coefficient of the CW potential (which includes contributions from various loops) using the latest WMAP3 data. Next, we use these results to further constrain $|K_i|$ and $\tilde{\lambda}_{1Z}$. We will finally comment on the implications of these constraints.

II. INFLATION WITH A COLEMAN-WEINBERG POTENTIAL

Let us now see under which conditions we can obtain a viable scenario for inflation with our model. As previously mentioned, the scalar field ϕ_Z receives various contributions to its potential, which will have the generic CW form [9]

$$V_0(\phi_Z^\dagger \phi_Z) = A\left(\phi_Z^\dagger \phi_Z\right)^2 \left(\log \frac{\phi_Z^\dagger \phi_Z}{v_Z^2} - \frac{1}{2}\right) + \frac{A v_Z^4}{2}. \quad (12)$$

After making the replacement (1) in (12), we obtain the potential for the real part σ_Z of ϕ_Z, which we want to be the inflaton field

$$V_0(\sigma_Z) = A(\sigma_Z + v_Z)^4 \left[\log \frac{(\sigma_Z + v_Z)^2}{v_Z^2} - \frac{1}{2}\right] + \frac{A v_Z^4}{2}. \quad (13)$$

This expression corresponds to the zero temperature limit. If we take into consideration finite-temperature effects, we should add a new term depending on temperature T that will give the following effective potential to σ_Z

$$V_{\rm eff}(\sigma_Z) = V_0(\sigma_Z) + \beta T^2 (\sigma_Z + v_Z)^2 \quad (14)$$

where β is a numerical constant. At high temperature, the field σ_Z is trapped at the $U(1)_A^{(Z)}$−symmetric minimum $\sigma_Z = -v_Z$. As the universe cools, for a sufficiently low temperature a new minimum appears at the $U(1)_A^{(Z)}$−symmetry breaking value $\sigma_Z = 0$ ($\langle \phi_Z \rangle = v_Z$). The critical temperature is the temperature at which the two minima become degenerate and is equal to $T_{\rm cr} = v_Z \sqrt{A/\beta}\ e^{-1/4}$. The universe cools further with the field σ_Z being trapped at the false vacuum and inflation starts when the false vacuum energy of σ_Z becomes dominant. Nevertheless, when the universe reaches the Hawking temperature

$$T_H = \frac{H}{2\pi} \simeq \frac{1}{2\pi}\sqrt{\frac{8\pi}{3M_{\rm P}^2} V_0(-v_Z)} = \frac{\sqrt{A}\, v_Z^2}{\sqrt{3\pi}\, M_{\rm P}} \quad (15)$$

a first-order phase transition occurs and σ_Z may start its slow-rolling towards the true minimum of the potential. In (15), $M_{\rm P} \simeq 1.22 \times 10^{19}$ GeV is the Planck-mass, H is the Hubble parameter at that epoch and we supposed that the energy density of σ_Z is the dominant one. Observable inflation occurs just after the false vacuum is destabilized and the inflaton slowly rolls down the potential. The evolution of σ_Z can be described classically.

Next, let us find the values for A and v_Z that are needed in order to obtain a viable model for inflation, compatible with observational data. The main constraints come from the combined observations of the Cosmic Microwave Background (CMB) and the Large Scale Structure (LSS) of the universe, which indicate the range of values for the spectral index $n_{\rm s}$, the tensor-to-scalar ratio r and, perhaps, evidence for a running in the spectral index. We also consider the constraint on the amplitude of the curvature perturbations, ${\mathcal{P}_\mathcal{R}}^{1/2}$, with the assumption that they were produced by quantum fluctuations of the inflaton field when the present large scales of the universe left the horizon during inflation. Finally, the number of e-folds of inflation produced between that epoch and the end of the inflationary stage should be large enough in order to solve the horizon and the flatness problems.

In our scenario, we have a theoretically motivated mechanism for generating a lepton asymmetry which then translates into a baryon asymmetry compatible with observations [3], [5]. Later on in this paper we will treat this aspect in more detail. For now, it is sufficient to say that after inflation, the σ_Z field starts to oscillate and to reheat the universe, mainly by decaying into two $\psi_i^{(Z)}$ fermions of masses given by (6), so that we want the inflaton to have sufficient mass as to decay into the two fermions. This is another condition to be considered for obtaining the adequate values for the parameters of our model.

Let us list the main constraints to be imposed on our model:

- the spectral index

$$n_{\rm s} \simeq 1 - 6\epsilon + 2\eta \quad (16)$$

  should be in the range $n_{\rm s} = 0.95 \pm 0.02$, where $\epsilon = \frac{M_{\rm P}^2}{16\pi}\left(\frac{V'}{V}\right)^2$ and $\eta = \frac{M_{\rm P}^2}{8\pi}\frac{V''}{V}$ are the slow-roll parameters and a prime means σ_Z-derivative;

- the right number of e-folds of inflation between large scale horizon crossing and the end of inflation

$$N_0 = \int_{\sigma_{Z,end}}^{\sigma_{Z,0}} \frac{V}{V'} d\sigma_Z \quad (17)$$

  where $\sigma_{Z,0}$ is the value of the inflaton field at horizon crossing and $\sigma_{Z,end}$ its value at the end of inflation;

- the amplitude of the curvature perturbations generated by the inflaton, evaluated at $\sigma_{Z,0}$

$${\mathcal{P}_\mathcal{R}}^{1/2} = \sqrt{\frac{128\pi}{3}}\frac{V^{3/2}}{M_{\rm P}^3 |V'|}\Big|_{\sigma_{Z,0}} \quad (18)$$

should have the WMAP3 value ${\mathcal{P}_{\mathcal{R}}}^{1/2} \simeq 4.7\times10^{-5}$;

- the inflaton mass $m_{\sigma_Z} = \sqrt{8A}\, v_Z$ should be at least 400 GeV or so in order for $\psi_i^{(Z)}$'s to be produced by the inflaton decay.

In our analysis, the parameters are functions of the inflaton field σ_Z and are evaluated when the present horizon scales left the inflationary horizon.

We have performed a complete numerical study. Imposing the requirements we have mentioned, we are interested in the lowest possible scale for inflation in our model. The scale is lowest for $v_Z \simeq 3\times10^9$ GeV. With this value, we obtain $A \simeq 3\times10^{-15}$ and $m_{\sigma_Z} \simeq 450$ GeV which, as commented above, is sufficient to produce two $\psi_i^{(Z)}$ fermions of masses $\mathcal{O}(200)$ GeV. We also obtain $n_{\rm s} = 0.923$ for the spectral index, not far away from the observed range, and $N \simeq 38$ e-folds of inflation between the present large scales horizon crossing until the end of inflation. The inflation scale is $V_0^{1/4} \equiv V(\sigma_{Z,0})^{1/4} \simeq 6\times10^5$ GeV. Low-scale inflation is interesting because it might be proved more easily in particle physics experiments.

Our model also satisfies the constraints for values of the parameters leading to higher values than $V_0^{1/4} \simeq 6\times10^5$ GeV. In Fig. 1 we show the dependence of the spectral index *vs* the energy scale of inflation. We see that the values of $n_{\rm s}$ are within 95% confidence level even at scales as low as 6×10^5 GeV, and increases with increasing inflationary scale. The graphic displayed in Fig. 1 was obtained with the assumption of instant reheating and a standard thermal history of the universe [11]. In Section III we present a more detailed analysis on the reheating mechanism. Here we just mention that the reheating temperature in our model is smaller than $V_0^{1/4}$, in which case the values of $n_{\rm s}$ displayed in Fig.1 shift to smaller values.

From now on we will stick to the lowest possible example $v_Z = 3\times10^9$ GeV ($V_0^{1/4} \sim 6\times10^5$ GeV) and examine the consequences. We should stress that the values of A does not vary drastically when we raise $V_0^{1/4}$ and we can safely consider it constant, with the value $A \simeq 3\times10^{-15}$. We then study some of the consequences that arise when adopting this value for A.

As stated before, one can have one-loop contributions to the parameter A coming from loops containing a) fermions $\psi_i^{(Z)}$, b) the messenger field $\tilde{\varphi}_1^{(Z)}$, and c) the inflaton. The fermion loop contribution, of order $-|K_i|^4/16\pi^2$, can be estimated for the values of the parameters chosen in the previous numerical example. From (7) we get for $v_Z = 3\times10^9$ GeV

$$|K_i| \lesssim 6.7\times10^{-8} \tag{19}$$

which will then translate into the following contribution to the A parameter in the CW potential

$$A_\psi \approx -|K_i|^4/16\pi^2 \sim 10^{-31} \tag{20}$$

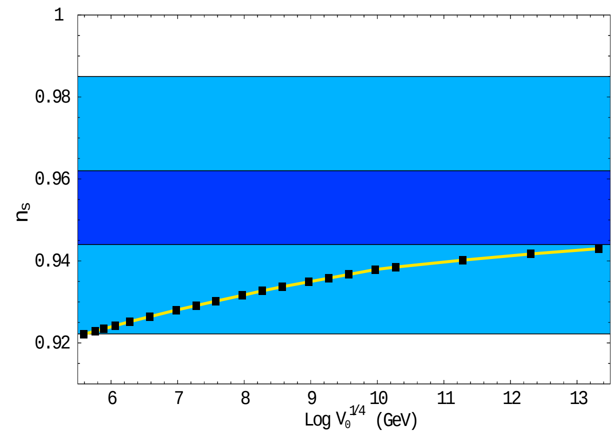

FIG. 1: The spectral index $n_{\rm s}$ as a function of the logarithm of the scale of inflation, $\log V_0^{1/4}$, compared with WMAP3 [1] range for $n_{\rm s}$ (at 68% and 95% confidence levels)

obviously being too small to be considered as contributing to it. Thus, fermion loops are completely negligible.

Next, we want to estimate what the contribution of the messenger field $\tilde{\varphi}_1^{(Z)}$ is. The leptogenesis bound (11) for $v_Z = 3\times10^9$ GeV becomes

$$\tilde{\lambda}_{1Z} \lesssim 10^{-13} \tag{21}$$

which gives the following contribution to the potential

$$A_{\tilde{\varphi}} \approx \tilde{\lambda}_{1Z}^2/16\pi^2 \sim 8\times10^{-29} \tag{22}$$

also being too small compared to the value $A = 3\times10^{-15}$. This means that the main contribution should come from σ_Z self-coupling λ. The necessary value of the λ coupling can be estimated by comparing its contribution $A_\sigma \approx \lambda^2/16\pi^2$ with A

$$A \simeq A_\sigma \approx \lambda^2/16\pi^2 \simeq 3\times10^{-15} \tag{23}$$

from which we obtain the constraint on λ

$$\lambda \simeq 6.9\times10^{-7}. \tag{24}$$

To end the discussion regarding inflation in our model, we would like to add that in our numerical study we obtained a small value for the running of the power spectrum, $\alpha \equiv \frac{{\rm d}n_{\rm s}}{{\rm d}\ln k} \simeq -0.002$. Other parameter that might be of interest is the tensor-to-scalar ratio r, which is defined usually as

$$r = \frac{\mathcal{P}_T}{\mathcal{P}_R} \tag{25}$$

where $\mathcal{P}_T$ and $\mathcal{P}_R$ are the power spectra for tensor and scalar perturbations, respectively. In the slow-roll regime of inflation, r can be expressed in terms of the slow-roll parameters and, at first order, $r = 16\epsilon$, where ϵ has to be evaluated at horizon crossing. With the values used in our previous numerical example, we obtain a very small tensor-to-scalar ratio $r \sim 10^{-43}$, making the quest for gravitational wave detection from the inflationary epoch hopeless.

It is amusing to note that the value of the σ_Z self-coupling $\lambda \sim O(10^{-7})$ that is consistent with the data is of the same order as the constraint on the Yukawa coupling $|K_i|$ coming from the CDM scenario of [3].

III. REHEATING

One of the most important questions of any inflationary scenario is the following: How do SM particles get generated at the end of inflation? In a generic inflationary model, it is not easy to answer this question since a generic inflaton is usually not coupled, either directly or indirectly, to SM particles. Although our inflaton is a SM-singlet field, we will show that its decay and the subsequent thermalization of the decay products can generate SM particles. In what follows, we will assume that the inflaton decays perturbatively as with the "old" reheating scenario and study its consequences. The interesting question of whether or not it can decay through the parametric resonance mechanism [12] of "preheating" scenarios is beyond the scope of this paper and will be dealt with separately elsewhere.

At the end of inflation, the inflaton will rapidly roll down its potential to the true minimum. The reheating (or, equivalently, the damping of the inflaton oscillation) occurs via the decay

$$\sigma_Z \to \psi^{(Z)} + \bar{\psi}^{(Z)}. \tag{26}$$

The width of the decay (26) is given by

$$\Gamma(\sigma_Z \to 2\psi^{(Z)}) \simeq 9\left(\frac{m_\psi}{v_Z}\right)^2 \frac{m_\sigma}{8\pi}\beta \tag{27}$$

where $\beta = (1-4(m_\psi/m_\sigma)^2)^{3/2}$ and we remember that $m_\sigma = \sqrt{8A}\, v_Z$. To estimate the reheating temperature caused by the process (26) after the end of inflation, we write

$$\Gamma(\sigma_Z \to 2\psi^{(Z)}) \sim H_{\rm rh} \tag{28}$$

where $H_{\rm rh} \sim 1.66 T_{\rm rh}^2/M_{\rm P}$ is the Hubble parameter at the reheating temperature $T_{\rm rh}$. By combining Eqs. (27) and (28) we obtain the dependence of the reheating temperature $T_{\rm rh}$ on v_Z

$$T_{\rm rh} \simeq 1.3\times 10^8 \left(\frac{v_Z}{\rm GeV}\right)^{-1/2}. \tag{29}$$

We see that the reheating temperature is a decreasing function of v_Z. This will set an upper bound on v_Z, because $T_{\rm rh}$ should be larger than twice the mass of $\psi^{(Z)}$ in order for the reheating mechanism to work, i.e.

$$T_{\rm rh} > 2m_\psi \sim 400 \text{ GeV} \tag{30}$$

which combined with (29) gives

$$v_Z < 10^{11} \text{ GeV}. \tag{31}$$

This upper limit restrict us to a low-scale inflation range, 6×10^5 GeV $\lesssim V_0^{1/4} \lesssim 2\times 10^7$ GeV, and then great part of Fig. 1 will be excluded, unless some other reheating mechanism is invoked. The spectral index values as a function of the logarithm of the scale of inflation, in the allowed range, is shown in Fig. 2. Notice that the values of $n_{\rm s}$ are a bit smaller now than in the case of instant reheating, but still marginally compatible with the WMAP3 value for $n_{\rm s}$.

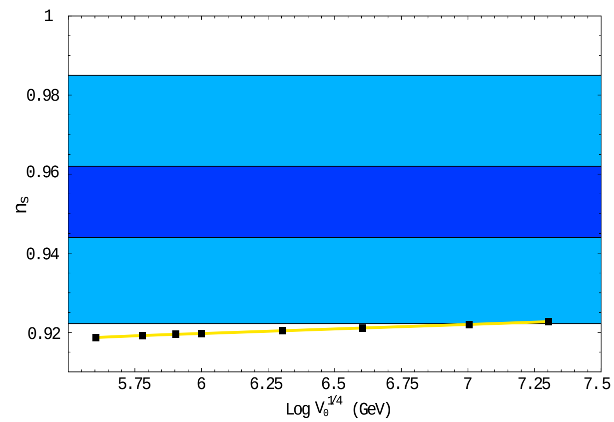

FIG. 2: The spectral index $n_{\rm s}$ as a function of the logarithm of the inflationary scale $V_0^{1/4}$, compared with WMAP3 [1] range for $n_{\rm s}$ (at 68% and 95% confidence levels), in the range allowed after imposing constraints coming from the reheating mechanism

Let us focus now on the mechanism by which SM particles are produced. For the sake of clarity in the following discussion, we will denote the QCD gluons by $\tilde{g}$ and the $SU(2)_Z$ "gluons" by $\mathbf{G}$. The chain of reactions which finally leads to the SM particles can be seen as follows:

$$\psi^{(Z)} + \bar{\psi}^{(Z)} \to \mathbf{G}\,\mathbf{G} \to \tilde{\varphi}_1^{(Z)}\,\tilde{\varphi}_1^{(Z)} \to W\,W,\, Z\,Z \to q\,\bar{q},\, l\,\bar{l}, \tag{32}$$

and

$$q\,\bar{q} \to \tilde{g}\tilde{g}\,. \tag{33}$$

We end up with a thermal bath of SM and $SU(2)_Z$ particles. This thermalization is possible because of the simple fact that $\tilde{\varphi}_1^{(Z)}$ carries *both SM and* $SU(2)_Z$ *quantum numbers*. Another important point concerns the various reactions rate in (32,33). The corresponding amplitudes are proportional to $O(g^2)$, where g stands for either the $SU(2)_Z$ coupling or a typical SM coupling at an energy above the electroweak scale. From [3] and [4], it can be seen that the various gauge couplings are of the same order of magnitude for a large range of energy, from a typical GUT scale down to the electroweak scale. One can safely conclude that the various reaction rates are comparable in magnitudes and the thermalization process shown above is truly effective. In principle, the messenger field also couples to $\psi^{(Z)}$ and a SM lepton, as shown in [3], but this is irrelevant in the thermalization process because the corresponding Yukawa couplings are too small.

It is remarkable to notice also that, because of the quantum numbers of the messenger field, the decay of $\tilde{\varphi}_1^{(Z)}$ into a SM lepton and $\psi^{(Z)}$ can generate a net SM lepton number which is subsequently transmogrified into a net baryon number through the electroweak sphaleron process as shown in [5]. In other words, the crucial presence of the messenger field $\tilde{\varphi}_1^{(Z)}$ facilitates both the generation of SM particles through thermalization and the subsequent leptogenesis through its decay.

IV. CONCLUSIONS

In this paper, we show that the model presented in [3], which explains dark matter and dark energy, also provides a mechanism for inflation in the early universe. We find that it is conceivable to have a low-scale inflation.

The complete model contains a new gauge group $SU(2)_Z$ which grows strong at a scale $\Lambda \sim 3 \times 10^{-3}$ eV, with the gauge coupling at GUT scale comparable to the SM couplings at the same scale. In addition to the SM particles, the model contains new particles: $\psi^{(Z)}_{(L,R),i}(i = 1, 2)$ which are $SU(2)_Z$ triplets and SM singlets; $\tilde{\varphi}^{(Z)}_{1,2}$ which are the so-called messenger fields and carry charges of both $SU(2)_Z$ and SM groups; and ϕ_Z, which is a singlet complex scalar field. The model also contains a new global symmetry $U(1)^{(Z)}_A$, which is spontaneously broken by the v.e.v. of the scalar field, $\langle\phi_Z\rangle = v_Z$, and also explicitly broken at the scale $\Lambda_Z \ll v_Z$ by the $SU(2)_Z$ gauge anomaly. The real part of the complex scalar field, namely σ_Z, is identified with the inflaton field. We considered a CW-type of potential for σ_Z and obtained the constraints on the parameters of the model in order to have a right description of inflation. The angular part of the complex scalar field, namely a_Z, acquires a small mass due to the explicit breaking of $SU(2)_Z$ and is trapped in a false vacuum, being responsible for the dark energy of the universe. The new particles $\psi^{(Z)}_{(L,R),i}$, with masses of order 200 GeV, explain the dark matter. They are produced at reheating, by the decay of the inflaton. For values of the $SU(2)_Z$ breaking scale $v_Z \sim 3 \times 10^9$ GeV, we obtain a low-scale model of inflation, namely a scale of $\sim 6 \times 10^5$ GeV. Notice that, in order to have a realistic reheating mechanism, this "low-scale" is also bounded from above by $\sim 2\times10^7$ GeV as we have discussed in the last section. Because of this fact, our model is a bona-fide "low-scale" inflationary scenario. It is an exciting possibility because the model might be indirectly probed at future LHC experiments.

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[1] D. N. Spergel *et al.*, arXiv:astro-ph/0603449.
[2] P. Q. Hung, "A new strong interaction sector as the origin for the dark energy and dark arXiv:hep-ph/0504060.
[3] P. Q. Hung, Nucl. Phys. B **747**, 55 (2006) [arXiv:hep-ph/0512282].
[4] P. Q. Hung and Paola Mosconi, arXiv:hep-ph/0611001
[5] P. Q. Hung, arXiv:hep-ph/0604063.
[6] P. Q. Hung, in preparation
[7] E. Masso and G. Zsembinszki, JCAP **0602**, 012 (2006) [arXiv:astro-ph/0602166].
[8] B. Spokoiny, Phys. Lett. B **315**, 40 (1993) [arXiv:gr-qc/9306008]; P. J. E. Peebles and A. Vilenkin, Phys. Rev. D **59**, 063505 (1999) [arXiv:astro-ph/9810509]; M. Peloso and F. Rosati, JHEP **9912**, 026 (1999) [arXiv:hep-ph/9908271]; K. Dimopoulos and J. W. F. Valle, Astropart. Phys. **18**, 287 (2002) [arXiv:astro-ph/0111417].
[9] S. R. Coleman and E. Weinberg, Phys. Rev. D **7**, 1888 (1973).
[10] Q. Shafi and V. N. Senoguz, Phys. Rev. D **73**, 127301 (2006) [arXiv:astro-ph/0603830].
[11] for reviews see A. R. Liddle and D. H. Lyth, Phys. Rept. **231**, 1 (1993) [arXiv:astro-ph/9303019]; D. H. Lyth and A. Riotto, Phys. Rept. **314**, 1 (1999) [arXiv:hep-ph/9807278].
[12] L. Kofman, A. D. Linde and A. A. Starobinsky, Phys. Rev. Lett. **73**, 3195 (1994) [arXiv:hep-th/9405187]; Y. Shtanov, J. H. Traschen and R. H. Brandenberger, Phys. Rev. D **51**, 5438 (1995) [arXiv:hep-ph/9407247].